β decay and shape isomerism in ^{74}Kr

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We study the properties of ^{74}Kr, and particularly the Gamow Teller strength distribution, using a deformed selfconsistent HF+RPA method with Skyrme type interactions. Results are presented for two density-dependent effective two-body interactions, including the dependence on deformation of the HF energy that exhibits two minima at close energies and distant deformations, one prolate and one oblate. We study the role of deformation, residual interaction, pairing and RPA correlations on the Gamow Teller strength distribution. Results on moments of inertia and gyromagnetic factors, as well as on $E0$ and $M1$ transitions are also presented.

I. INTRODUCTION

The study of exotic nuclei, characterized by nonoptimal N/Z ratios, has attracted an increasing attention during the last decades from both experimental and theoretical points of view [1,2]. Nevertheless, it is still true that most of the knowledge we have at present about nuclear systems corresponds to the nuclei inside the valley of stability, where the balance between protons and neutrons make these systems stable against particle emission. There is still much to be learned about systems out of this region of stability. Many complementary experimental projects and new facilities involving radioactive nuclear beams are being carried out or planned at CERN, GANIL, Grenoble-Munich, GSI, Louvain-la-Neuve, RIKEN... to extend this limited knowledge.

Information on these nuclei is of prime importance to understand different problems in nuclear physics related to the way nuclei respond as we vary their masses and isospin parameters to extreme values of the N/Z ratio. One can also learn about specific components of the N-N force and the modification of this force in the nuclear medium. Since the isospin dependence of the N-N interaction is largely unknown, the structure of single particle states, collective modes and behaviour of global nuclear properties is very uncertain in nuclei with unusual N/Z ratios. The theoretical techniques and parameters are optimized to reproduce the properties of known atomic masses and it is not clear how accurate is the theoretical extrapolation to regions far from stability. The study of nuclei far from stability provides then new basic pieces of information on nuclear structure.

Knowledge of the decay properties of radioactive nuclei is also essential for nuclear astrophysics [3], where one tries to understand the production of energy and the synthesis of elements in stars during stellar events. In that respect, reactions with radioactive partners become very important, specially during violent stellar events with extreme conditions of density and temperature because the average time between successive nuclear reactions is much shorter than the average decay time of the radioactive nucleus. Therefore, these radioactive nuclei do not have time to decay before participating in new nuclear reactions. In particular, the decay properties of proton rich nuclei are fundamental to understand the r-p process (rapid proton capture nucleosynthesis) that involves multiple proton capture reactions on proton rich nuclei. The microscopic structure of exotic nuclei near the drip lines is therefore a topic of great interest in nuclear structure as well as in astrophysics.

The neutron deficient side of the valley of stability is particularly interesting because in this region a large fraction of the Gamow Teller (GT) strength is accessible in β^+ decay [4]. Usually kinematics limits β decay to a narrow energy window that in general misses the main part of the giant GT resonance, but this is not the case in proton rich nuclei. Therefore, proton rich nuclei provide a good opportunity to test the GT strength measured by other less direct means such as (p, n) reactions, where the GT strength is extracted by extrapolating the empirical proportionality between the zero degree cross section and the β decay strength to low-lying GT states [5].

All of these features, together with the ongoing experimental effort to measure GT strengths in proton rich medium mass nuclei [6], make this region the focus of this paper.

It is clear that, at some point, extensions of the standard approaches will be necessary to cover particular aspects that, although not important in stable nuclei, may play a role in describing the unstable ones. Issues such as the role of the continuum [7] in weakly bound systems or the neutron-proton pairing [8] in $N \approx Z$ nuclei should be considered. It is nevertheless important to know whether the most powerful tools, designed to reproduce the properties of nuclei within the valley of stability, are still valid when one approaches the drip lines. One should know the limits of applicability

of the existing interactions and methods. In our view this has not been sufficiently explored yet, particularly in what concerns the GT strength distribution.

With this view in mind we study in this paper the β^+ decay of ^{74}Kr, which is a medium mass proton rich nucleus of recent experimental [9] and theoretical interest. This nucleus is also representative of a mass region that exhibits a wide variety of shapes including large quadrupole deformations ($\beta_2 \simeq 0.4$) and shape coexistence. Our method consists of a selfconsistent formalism based on a deformed Hartree-Fock (HF) mean field obtained with a Skyrme interaction including pairing correlations in the BCS approximation. We add to this mean field a spin-isospin residual interaction with a coupling strength derived by integrating over the nuclear volume the Landau-Migdal force, obtained from the same Skyrme interaction. The residual force is therefore consistent with the mean field. The equations of motion are solved in the proton-neutron quasiparticle random phase approximation (QRPA) [10]. Most β decay calculations in the past used as input in one way or another experimental single particle levels, but the extrapolation of these empirical information to exotic nuclei is at least questionable and the need for selfconsistent calculations in unstable nuclei has been emphasized (see for instance [11]). There are also selfconsistent mean field studies [7], as well as investigations [12] of the dependence on deformation of the GT strength distributions in Tamm Dancoff approximation (TDA), but to our knowledge this is the first time that the selfconsistent deformed QRPA formalism described here is used to calculate GT strengths. A similar formalism for spin $M1$ strength distributions has been successfully tested [13,14] against experimental data in the rare earths and actinide nuclei.

The paper is organized as follows: In Section 2 we introduce the theoretical framework used to describe the β decay. We first introduce the selfconsistent method for the mean field and residual interactions and then we obtain the β strengths in bare two quasiparticle (qp) approximation, in TDA, and in QRPA. In Section 3 we study the ground state properties and the low-lying structure of ^{74}Kr and compare our calculations to experiment. In Section 4 we discuss the results on Gamow Teller strength distributions and sum rules obtained in our calculations studying the sensitivity of these results to the various ingredients of our approach such as the effective two-body interaction, the pairing interaction, the nuclear shape, and the RPA correlations. We also show results on $M1$ transitions. Section 5 contains half-life results in various approaches. Section 6 contains the main conclusions of this work.

II. BRIEF DESCRIPTION OF THEORY AND DETAILS OF CALCULATIONS

In this section we summarize the theory involved in the microscopic calculations presented in Sections 3 and 4. Our method is selfconsistent in the sense that, as described in Subsection A, with a given two-body density-dependent effective interaction we derive: a) the selfconsistent mean field generating the single particle energies, wave functions and occupations of the ground state, b) the particle-hole Landau-Migdal interaction averaged over the nuclear volume generating the QRPA modes. Thus, the same two-body interaction is used to derive the QRPA excitations consistently with the quasiparticle basis and ground state, as described in Subsection B.

A. Mean field calculations and residual interactions

It is well known that the density-dependent HF approximation gives a very good description of ground-state properties for both spherical and deformed nuclei [15] and it is at present the most reliable mean field description. We consider in this paper two different Skyrme forces. On the one hand we use the Skyrme force Sk3 [16] because it is the most extensively used Skyrme force and we consider it as a reference. We use Sk3 in its density dependent two-body version that has better spin-isospin properties than the three-body one [13,14]. On the other hand we use the SG2 force [17] of Van Giai and Sagawa. The two forces were designed to fit ground state properties of spherical nuclei and nuclear matter properties but, in addition, the force SG2 gives a good description of Gamow Teller excitations in spherical nuclei [17]. Recently [13], we applied these two forces and a similar method to make a rather extensive study of isoscalar and isovector spin M1 excitations in deformed nuclei obtaining a good description of the available data, particularly with SG2. This gives us confidence on the predictive power of the method and on the reliability of the SG2 force when considering isovector spin properties of deformed nuclei. The parameters of these two interactions are given in Table 1, and the corresponding HF energy density functional for an even-even nucleus has the form

$$
\mathcal{E}\left(\mathbf{r}\right) = \sum_{st} \rho_{st} \sum_{s't'} \left\{ \frac{1}{2} t_0 \rho_{s't'} \left[1 - \delta_{ss'}\delta_{tt'} + x_0 \left(\delta_{ss'} - \delta_{tt'} \right) \right] \right.
$$
$$
\left. + \frac{1}{4} t_2 \left(\tau_{s't'} + \frac{1}{4} \nabla^2 \rho_{s't'} \right) \left[1 + \delta_{ss'}\delta_{tt'} + x_2 \left(\delta_{ss'} + \delta_{tt'} \right) \right] \right.
$$

$$+\frac{1}{16}t_1\left(4\tau_{s't'}-3\nabla^2\rho_{s't'}\right)\left[1-\delta_{ss'}\delta_{tt'}+x_1\left(\delta_{ss'}-\delta_{tt'}\right)\right]$$

$$+\frac{1}{12}t_3\rho^\alpha\rho_{s't'}\left[1-\delta_{ss'}\delta_{tt'}+x_3\left(\delta_{ss'}-\delta_{tt'}\right)\right]$$

$$+\frac{i}{2}W_0\boldsymbol{\nabla}\cdot\mathbf{J}_{s't'}\left(1+\delta_{tt'}\right)\Big\}+\mathcal{E}_C\left(\mathbf{r}\right) \tag{1}$$

with $\mathcal{E}_C$ the Coulomb energy density

$$\mathcal{E}_C\left(\mathbf{r}\right)=e^2\frac{1}{2}\int d\mathbf{r}'\frac{\rho_p\left(\mathbf{r}'\right)\rho_p\left(\mathbf{r}\right)}{|\mathbf{r}-\mathbf{r}'|}-\frac{3}{4}e^2\rho_p\left(\mathbf{r}\right)\left[\frac{3}{\pi}\rho_p\left(\mathbf{r}\right)\right]^{1/3} \tag{2}$$

The spin-isospin (st) components of the nucleon, kinetic energy, and magnetization densities are

$$\rho_{st}\left(\mathbf{r}\right)=\sum_i v_i^2\left|\phi_i\left(\mathbf{r},s,t\right)\right|^2 \tag{3}$$

$$\tau_{st}\left(\mathbf{r}\right)=\sum_i v_i^2\left|\boldsymbol{\nabla}\phi_i\left(\mathbf{r},s,t\right)\right|^2 \tag{4}$$

$$\mathbf{J}_{st}\left(\mathbf{r}\right)=\sum_{i,s'}v_i^2\phi_i^*\left(\mathbf{r},s',t\right)\left(-i\boldsymbol{\nabla}\times\boldsymbol{\sigma}\right)\phi_i\left(\mathbf{r},s,t\right) \tag{5}$$

with

$$\rho_t=\sum_s\rho_{st}\ , \tag{6}$$

$$\rho=\sum_{t=p,n}\rho_t \tag{7}$$

and similarly for τ and $\mathbf{J}$.

In our calculations time-reversal and axial symmetry are assumed. The single-particle wave functions are expanded in terms of the eigenstates of an axially symmetric harmonic oscillator in cylindrical coordinates. The single-particle states $|i\rangle$ and their time reversed $|\bar{i}\rangle$ are characterized by the eigenvalues Ω of J_z, parity π_i, and energy ϵ_i

$$|i\rangle=\sum_N\frac{(-1)^N+\pi_i}{2}\sum_{n_r,n_z,\Lambda\geq 0,\Sigma}C^i_{Nn_rn_z\Lambda\Sigma}\left|Nn_rn_z\Lambda\Sigma\right\rangle \tag{8}$$

with $\Omega_i=\Lambda+\Sigma\geq\frac{1}{2}$, and

$$|\bar{i}\rangle=\sum_N\frac{(-1)^N+\pi_i}{2}\sum_{n_r,n_z,\Lambda\geq 0,\Sigma}C^i_{Nn_rn_z\Lambda\Sigma}(-1)^{\frac{1}{2}-\Sigma}\left|Nn_rn_z-\Lambda-\Sigma\right\rangle \tag{9}$$

with $\Omega_{\bar{i}}=-\Omega_i=-\Lambda-\Sigma\leq-\frac{1}{2}$. For each N the sum over $n_r,n_z,\Lambda\geq 0$ is extended to the quantum numbers satisfying $2n_r+n_z+\Lambda=N$. The sum over N goes from $N=0$ to $N=10$ in our calculations.

The single-particle energies ϵ_i and wave functions are obtained from the HF equations

$$\frac{\delta\mathcal{E}}{\delta\phi_i^*}=\epsilon_i\phi_i \tag{10}$$

We use the McMaster code that is based in the formalism developed in Ref. [18] as described in Ref. [19]. The method includes pairing between like nucleons in the BCS approximation with fixed gap parameters for protons, Δ_p, and neutrons, Δ_n. The number equation in the neutron sector reads

$$2\sum_i v_i^2=N \tag{11}$$

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where v_i^2 are the occupation probabilities

$$v_i^2 = \frac{1}{2}\left[1 - \frac{\epsilon_i - \lambda_n}{E_i}\right] \; ; \; u_i^2 = 1 - v_i^2 \tag{12}$$

in terms of the quasiparticle energies

$$E_i = \sqrt{(\epsilon_i - \lambda_n)^2 + \Delta_n^2} \tag{13}$$

Eq. (11) is solved iteratively at the end of each Hartree-Fock iteration to determine the Fermi level λ_n. Similar equations are used to determine the Fermi level and occupation probabilities for protons by changing N into Z, Δ_n into Δ_p, and λ_n into λ_p.

The fixed gap parameters are determined phenomenologically from the odd-even mass differences through a symmetric five term formula involving the experimental binding energies [20]:

$$\Delta_n = \frac{1}{8}\left[B\left(N - 2, Z\right) - 4B\left(N - 1, Z\right) + 6B\left(N, Z\right)\right.$$
$$\left. -4B\left(N + 1, Z\right) + B\left(N + 2, Z\right)\right] \tag{14}$$

A similar expression is found for the proton gap Δ_p by changing N by Z and vice versa. For ^{74}Kr we obtain $\Delta_n = \Delta_p = 1.5$ MeV.

The HF theory gives a single solution which is the Slater determinant of lowest energy. To allow for shape coexistence one has to extend the theory to a constrained HF theory [21]. Minimization of the HF energy under the constraint of holding the nuclear deformation fixed is carried out over a range of deformations. When more than one local minimum occurs for the total energy as a function of deformation, shape coexistence results. The energy surfaces as a function of deformation that we will discuss in Section 3, are obtained by this procedure including a quadratic quadrupole constraint [21].

Following Bertsch and Tsai [22] the particle-hole interaction consistent with the HF mean field can be obtained as

$$V_{ph} = \frac{1}{16}\sum_{sts't'}\left[1 + (-1)^{s-s'}\boldsymbol{\sigma_1}\cdot\boldsymbol{\sigma_2}\right]\left[1 + (-1)^{t-t'}\boldsymbol{\tau_1}\cdot\boldsymbol{\tau_2}\right]\frac{\delta^2\mathcal{E}}{\delta\rho_{st}\left(\mathbf{r}_1\right)\delta\rho_{s't'}\left(\mathbf{r}_2\right)} \tag{15}$$

This gives a local interaction that can be put in the Landau-Migdal form [23]. For the study of β decay the relevant residual interactions are the isospin contact forces generating the allowed Fermi transitions ($\Delta L = 0, \Delta S = 0, \Delta I^\pi = 0^+$)

$$V_F\left(12\right) = \chi_F\left(t_1^+ t_2^- + t_1^- t_2^+\right) \tag{16}$$

and the spin-isospin contact forces generating the allowed Gamow Teller transitions ($\Delta L = 0, \Delta S = 1, \Delta I^\pi = 1^+$)

$$V_{GT}\left(12\right) = \chi_{GT}\,\boldsymbol{\sigma_1}\cdot\boldsymbol{\sigma_2}\left(t_1^+ t_2^- + t_1^- t_2^+\right) \tag{17}$$

where we use the convention $t^+\left|p\right\rangle = \left|n\right\rangle,\ t^-\left|n\right\rangle = \left|p\right\rangle$. The latter ($V_{GT}$) is the charge changing component of the spin-spin interaction H_{SS}

$$H_{SS} = \frac{K_S}{4A}\left[(1 + q)\,\mathbf{s}_1\cdot\mathbf{s}_2 + (1 - q)\,\mathbf{s}_1\cdot\mathbf{s}_2\,\boldsymbol{\tau_1}\cdot\boldsymbol{\tau_2}\right] \tag{18}$$

considered in Refs. [13,14] for the study of spin $M1$ excitations. Not allowed transitions ($\Delta L > 0$) produce strengths which are orders of magnitude smaller than the allowed ones ($\Delta L = 0$) and will not be considered in this work.

After functional differentiation we get from Eqs. (1) and (15-17), assuming symmetric uniform nuclear matter and averaging over the nuclear volume V

$$\chi_F = \frac{3}{4\pi R^3}\left(-\frac{1}{2}\right)\left\{t_0\left(1 + 2x_0\right) - \frac{1}{2}k_F^2\left[t_2\left(1 + 2x_2\right) - t_1\left(1 + 2x_1\right)\right] + \frac{1}{6}t_3\rho^\alpha\left(1 + 2x_3\right)\right\} \tag{19}$$

$$\chi_{GT} = \frac{3}{4\pi R^3}\left(-\frac{1}{2}\right)\left\{t_0 + \frac{1}{2}k_F^2\left(t_1 - t_2\right) + \frac{1}{6}t_3\rho^\alpha\right\} \tag{20}$$

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where R is the nuclear radius and k_F the Fermi momentum $k_F = \left(3\pi^2\rho/2\right)^{1/3}$. These coupling strengths are related to the familiar Landau-Migdal parameters F_0' and G_0' (see for instance [17]) by

$$\chi_F = \frac{2F_0'}{V N_0} \; , \; \chi_{GT} = \frac{2G_0'}{V N_0} \tag{21}$$

where $V = 4\pi R^3/3$ and $N_0 = \left(2m^\star k_F/\hbar^2\pi^2\right)$, with $m^\star$ the effective mass. For completeness the values of χ_F and χ_{GT} are also given in Table 1.

B. QRPA equations and β decay strengths

Let us first consider Gamow Teller excitations. Using the notation $\beta_K^\pm = \sigma_K t^\pm$ $(K = 0, \pm 1)$, we write Eq. (17) in the more traditional form

$$V_{GT} = 2\chi_{GT} \sum_K (-1)^K \; \beta_K^+ \beta_{-K}^- \tag{22}$$

where in second quantization we can write

$$\beta_K^+ = \sum_{np} \langle n \left| \sigma_K \right| p \rangle a_n^+ a_p^+ = \sum_{np} \langle n \left| \sigma_K \right| p \rangle \left\{ u_n v_p \alpha_n^+ \alpha_{\bar{p}}^+ \right.$$
$$\left. + v_n u_p \alpha_{\bar{n}} \alpha_p + u_n u_p \alpha_n^+ \alpha_p + v_n v_p \alpha_{\bar{p}}^+ \alpha_{\bar{n}} \right\} \tag{23}$$

and $\beta_{-K}^- = (-1)^K \left(\beta_K^+\right)^+$. We use the standard notation a^+ (a) for particle creation (annihilation) and α^+ (α) for quasiparticle creation (annihilation) operators. We write the proton-neutron QRPA phonon operator for Gamow Teller excitations in even-even nuclei

$$\Gamma_{\omega_K}^+ = \sum_{\gamma_K} \left[X_{\gamma_K}^{\omega_K} A_{\gamma_K}^+ - Y_{\gamma_K}^{\omega_K} A_{\bar{\gamma}_K} \right] \tag{24}$$

where $A_{\gamma_K}^+$, $A_{\bar{\gamma}_K}$ are two-quasiparticle operators of the form

$$A_{\gamma_K}^+ = \alpha_n^+ \alpha_{\bar{p}}^+ \tag{25}$$

$$A_{\bar{\gamma}_K} = \left(-\alpha_{\bar{n}}^+ \alpha_p^+\right)^+ = \alpha_{\bar{n}} \alpha_p \tag{26}$$

with $\Omega_n + \Omega_{\bar{p}} = K$, where n (p) stands for any neutron (proton) states $|i\rangle$ or $|\bar{i}\rangle$ and $\bar{n}$ $(\bar{p})$ is its time reverse. The QRPA equations for the even-even system are

$$\left\langle \phi_0 \left| A_{\gamma_K} \left[H, \Gamma_{\omega_K}^+\right] \right| \phi_0 \right\rangle = \omega_K \left\langle \phi_0 \left| A_{\gamma_K} \Gamma_{\omega_K}^+ \right| \phi_0 \right\rangle \tag{27}$$

$$\left\langle \phi_0 \left| \left[H, \Gamma_{\omega_K}^+\right] A_{\bar{\gamma}_K}^+ \right| \phi_0 \right\rangle = \omega_K \left\langle \phi_0 \left| \Gamma_{\omega_K}^+ A_{\bar{\gamma}_K}^+ \right| \phi_0 \right\rangle \tag{28}$$

with

$$H = \sum_n \alpha_n^+ \alpha_n E_n + \sum_p \alpha_p^+ \alpha_p E_p + V_{GT} - \left\langle \phi_0 \left| H \right| \phi_0 \right\rangle \tag{29}$$

$|\phi_0\rangle$ is the HF+BCS ground state satisfying $\alpha_t |\phi_0\rangle = 0$ for any t. From these equations the forward and backward amplitudes are obtained as

$$X_{\gamma_K}^{\omega_K} = \frac{2\chi_{GT}}{\omega_K - \mathcal{E}_{\gamma_K}} \left(a_{\gamma_K} \mathcal{M}_+^{\omega_K} + b_{\gamma_K} \mathcal{M}_-^{\omega_K}\right) \tag{30}$$

$$Y_{\gamma_K}^{\omega_K} = \frac{-2\chi_{GT}}{\omega_K + \mathcal{E}_{\gamma_K}} \left(b_{\gamma_K} \mathcal{M}_+^{\omega_K} + a_{\gamma_K} \mathcal{M}_-^{\omega_K}\right) \tag{31}$$

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with $\mathcal{E}_{\gamma_K} = E_n + E_p$ the two-quasiparticle excitation energy with E_i given by Eq. (13), $\mathcal{M}^{\omega_K}_{\pm}$ are

$$\mathcal{M}^{\omega_K}_{+} = \sum_{\gamma_K} \left(a_{\gamma_K} X^{\omega_K}_{\gamma_K} + b_{\gamma_K} Y^{\omega_K}_{\gamma_K} \right) \tag{32}$$

$$\mathcal{M}^{\omega_K}_{-} = \sum_{\gamma_K} \left(b_{\gamma_K} X^{\omega_K}_{\gamma_K} + a_{\gamma_K} Y^{\omega_K}_{\gamma_K} \right) \tag{33}$$

with

$$a_{\gamma_K} = u_n v_p \Sigma^{np}_K \tag{34}$$

$$b_{\gamma_K} = v_n u_p \Sigma^{np}_K \tag{35}$$

and

$$\Sigma^{np}_K = \langle n \left| \sigma_K \right| p \rangle \tag{36}$$

Owing to the transformation properties under time reversal and hermitian conjugation it is easy to show that: i) the $K = -1$ and $K = 1$ modes are related to each other through time reversal and are degenerate; ii) the $K = 0$ modes are time odd and satisfy $X^{\omega_0}_{\tilde{\gamma}_0} = -X^{\omega_0}_{\gamma_0}$ and $Y^{\omega_0}_{\tilde{\gamma}_0} = -Y^{\omega_0}_{\gamma_0}$. Forward and backward amplitudes are real. This allows to simplify the calculations to include only neutron-proton pairs with $\Omega_n > 0$, i.e., the complete set of γ_K neutron proton pairs can be reduced to the basic sets of pairs:

- $i_0 = n_i p_i$ with $\Omega_{n_i} = \Omega_{p_i} \geq \frac{1}{2}$

- $i_{-1} = n_i p_i$ with $\Omega_{n_i} = \Omega_{p_i} - 1 \geq \frac{1}{2}$

- $i_{+1} = \begin{cases} n_i p_i & \text{with } \Omega_{n_i} = \Omega_{p_i} + 1 \geq \frac{3}{2} \\ n_i \bar{p}_i & \text{with } \Omega_{n_i} = \Omega_{p_i} = \frac{1}{2} \end{cases}$

The matrix elements are

$$\Sigma^{n_i p_i}_K = \sum_{N n_z \Lambda \Sigma} C^{n_i}_{N n_z \Lambda \Sigma + K} C^{p_i}_{N n_z \Lambda \Sigma} \left(2\Sigma \right) \sqrt{1 + |K|} \tag{37}$$

$$\Sigma^{n_i \bar{p}_i}_{K=1} = \sum_{N n_z} C^{n_i}_{N n_z 0 \frac{1}{2}} C^{p_i}_{N n_z 0 \frac{1}{2}} \left(-\sqrt{2} \right) \tag{38}$$

We therefore obtain the following dispersion relations for the $K = 0$ and $K = 1$ modes:

$$\left(\frac{1}{4\chi_{GT}} \right)^2 = \frac{1}{2\chi_{GT}} \sum_{i_0} \frac{\left(a^2_{i_0} + b^2_{i_0} \right)}{\omega^2_0 - \mathcal{E}^2_{i_0}} \mathcal{E}_{i_0} + \left(\sum_{i_0} a_{i_0} b_{i_0} \frac{2\mathcal{E}_{i_0}}{\omega^2_0 - \mathcal{E}^2_{i_0}} \right)^2$$
$$- \sum_{i_0} \left(\frac{a^2_{i_0}}{\omega_0 + \mathcal{E}_{i_0}} - \frac{b^2_{i_0}}{\omega_0 - \mathcal{E}_{i_0}} \right) \sum_{i_0} \left(\frac{b^2_{i_0}}{\omega_0 + \mathcal{E}_{i_0}} - \frac{a^2_{i_0}}{\omega_0 - \mathcal{E}_{i_0}} \right) \tag{39}$$

for $K = 0$, and

$$\left(\frac{1}{2\chi_{GT}} \right)^2 = \frac{1}{\chi_{GT}} \sum_{i_\rho (\rho = \pm 1)} \frac{\left(a^2_{i_\rho} + b^2_{i_\rho} \right)}{\omega^2_1 - \mathcal{E}^2_{i_\rho}} \mathcal{E}_{i_\rho} + \left(\sum_{i_\rho (\rho = \pm 1)} a_{i_\rho} b_{i_\rho} \frac{2\mathcal{E}_{i_\rho}}{\omega^2_1 - \mathcal{E}^2_{i_\rho}} \right)^2$$
$$- \sum_{i_\rho (\rho = \pm 1)} \left(\frac{a^2_{i_\rho}}{\omega_1 + \mathcal{E}_{i_\rho}} - \frac{b^2_{i_\rho}}{\omega_1 - \mathcal{E}_{i_\rho}} \right) \sum_{i_\rho (\rho = \pm 1)} \left(\frac{b^2_{i_\rho}}{\omega_1 + \mathcal{E}_{i_\rho}} - \frac{a^2_{i_\rho}}{\omega_1 - \mathcal{E}_{i_\rho}} \right) \tag{40}$$

for $K = 1$.

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The forward and backward amplitudes are then obtained from Eqs. (30) and (31) with the following normalization conditions

$$2\sum_{i_0}\left[\left(X_{i_0}^{\omega_0}\right)^2-\left(Y_{i_0}^{\omega_0}\right)^2\right]=1 \tag{41}$$

$$\sum_{i_\rho(\rho=\pm1)}\left(X_{i_\rho}^{\omega_1}\right)^2-\left(Y_{i_\rho}^{\omega_1}\right)^2=1 \tag{42}$$

Using the inverse transformation

$$A_{i_K}^+=\sum_{\omega_K}\left(X_{i_K}^{\omega_K}\Gamma_{\omega_K}^++Y_{i_K}^{\omega_K}\Gamma_{\bar{\omega}_K}\right), \tag{43}$$

where $\Gamma_{\bar{\omega}_K}$ stands for the time reverse of Γ_{ω_K}, and the RPA conditions $\Gamma_{\omega_K}\left|0\right\rangle=0$, $\left|\omega_K\right\rangle=\Gamma_{\omega_K}^+\left|0\right\rangle$, with $\left|0\right\rangle$ the QRPA correlated ground state, one gets the $\beta_K^\pm$ strengths

$$\left\langle\omega_K\left|\beta_K^\pm\right|0\right\rangle=\mathcal{M}_\pm^{\omega_K} \tag{44}$$

with

$$\mathcal{M}_+^{\omega_0}=2\sum_{i_0}\left(a_{i_0}X_{i_0}^{\omega_0}+b_{i_0}Y_{i_0}^{\omega_0}\right) \tag{45}$$

$$\mathcal{M}_+^{\omega_1}=\sum_{i_\rho(\rho=\pm1)}\left(a_{i_\rho}X_{i_\rho}^{\omega_1}+b_{i_\rho}Y_{i_\rho}^{\omega_1}\right) \tag{46}$$

$\mathcal{M}_-^{\omega_K}$ is obtained from $\mathcal{M}_+^{\omega_K}$ exchanging a_{i_K} and b_{i_K} (see Eq. (33)).

To finish this section we note that the Fermi transitions, although not shown here, can be handled in a similar way to the $K=0$ GT transitions with the important difference that the phonon operator for the Fermi mode is time even. Therefore, if we denote by ω_F the frequencies of the Fermi excitations one has that $X_{\bar{i}_0}^{\omega_F}=X_{i_0}^{\omega_F}$, $Y_{\bar{i}_0}^{\omega_F}=Y_{i_0}^{\omega_F}$ and we may consider only pairs of the type i_0 with the matrix elements a_{i_0}, b_{i_0} replaced by

$$a_{i_0}'=\left\langle n_i|p_i\right\rangle u_{n_i}v_{p_i} \tag{47}$$

$$b_{i_0}'=\left\langle n_i|p_i\right\rangle v_{n_i}u_{p_i} \tag{48}$$

With these replacements the dispersion relation in Eq. (39) can also be used to find the frequencies ω_F of the Fermi modes and the Fermi strengths can be obtained similarly to the $K=0$ GT strengths

$$\left\langle\omega_F\left|t^+\right|0\right\rangle=\mathcal{M}_+^{\omega_F}=2\sum_{i_0}\left(a_{i_0}'X_{i_0}^{\omega_F}+b_{i_0}'Y_{i_0}^{\omega_F}\right) \tag{49}$$

$$\left\langle\omega_F\left|t^-\right|0\right\rangle=\mathcal{M}_-^{\omega_F}=2\sum_{i_0}\left(a_{i_0}'Y_{i_0}^{\omega_F}+b_{i_0}'X_{i_0}^{\omega_F}\right) \tag{50}$$

with

$$X_{i_0}^{\omega_F}=\left(a_{i_0}'\mathcal{M}_+^{\omega_F}+b_{i_0}'\mathcal{M}_-^{\omega_F}\right)\frac{2\chi_F}{\omega_F-\mathcal{E}_{i_0}} \tag{51}$$

$$Y_{i_0}^{\omega_F}=-\left(b_{i_0}'\mathcal{M}_+^{\omega_F}+a_{i_0}'\mathcal{M}_-^{\omega_F}\right)\frac{2\chi_F}{\omega_F+\mathcal{E}_{i_0}} \tag{52}$$

and

$$2 \sum_{i_0} \left(\left| X_{i_0}^{\omega_F} \right|^2 - \left| Y_{i_0}^{\omega_F} \right|^2 \right) = 1$$

In Section 4 we show also for comparison the Gamow Teller strength distributions obtained in the TDA approximation, as well as the uncorrelated two-quasiparticle neutron proton GT strengths.

The equations for TDA are easily obtained from the above RPA equations by taking everywhere the limit $1/ \left(\omega_K + \mathcal{E}_{i_K} \right) \to 0$, that in particular kills the Y amplitudes $\left(Y_{i_K}^{\omega_K} = 0 \right)$ and simplifies the eigenvalue equations. The $\beta_K^{\pm}$ strengths are then given by

$$\left\langle \omega_K^{TDA} \left| \beta_K^{\pm} \right| \phi_0 \right\rangle = \mathcal{M}_{\pm}^{\omega_K^{TDA}} \tag{53}$$

with

$$\mathcal{M}_{+}^{\omega_0^{TDA}} = 2 \sum_{i_0} a_{i_0} X_{i_0}^{\omega_0^{TDA}} \tag{54}$$

$$\mathcal{M}_{-}^{\omega_0^{TDA}} = 2 \sum_{i_0} b_{i_0} X_{i_0}^{\omega_0^{TDA}} \tag{55}$$

$$2 \sum_{i_0} \left(X_{i_0}^{\omega_0^{TDA}} \right)^2 = 1$$

and corresponding relations for $K = 1$.

Finally, the uncorrelated two-quasiparticle neutron-proton Gamow Teller (Fermi) excitations are obtained in the limit $\chi_{GT} = 0$ $(\chi_F = 0)$. Obviously in this case the excitation energies are the bare two-quasiparticle energies $\omega_K^{2qp} = \mathcal{E}_{\gamma_K} = E_n + E_p$, and the $\beta^{\pm}$ strengths are

$$\left\langle \omega_K^{2qp} \left| \beta_K^{+} \right| \phi_0 \right\rangle = a_{\gamma_K} \tag{56}$$

$$\left\langle \omega_K^{2qp} \left| \beta_K^{-} \right| \phi_0 \right\rangle = b_{\gamma_K} \tag{57}$$

Similar relations hold for Fermi transitions with $K = 0$ and a_{γ_0}, b_{γ_0} replaced by a'_{γ_0}, b'_{γ_0} (47) and (48).

In the uncorrelated case it is straightforward to show that the Ikeda sum rule is fulfilled. For Fermi transitions one has

$$S_F^{-} - S_F^{+} = \sum_{\omega_F} \left| \left\langle \omega_F \left| t^{-} \right| \phi_0 \right\rangle \right|^2 - \sum_{\omega_F} \left| \left\langle \omega_F \left| t^{+} \right| \phi_0 \right\rangle \right|^2 = \sum_{\gamma_0} \left[\left(b'_{\gamma_0} \right)^2 - \left(a'_{\gamma_0} \right)^2 \right]$$

$$= 2 \sum_{i_0} \left[\left(b'_{i_0} \right)^2 - \left(a'_{i_0} \right)^2 \right] = 2 \sum_{n_i, p_i} \left| \left\langle n_i | p_i \right\rangle \right|^2 \left(v_{n_i}^2 - v_{p_i}^2 \right) = N - Z \tag{58}$$

For Gamow Teller transitions one has

$$S_{GT}^{-} - S_{GT}^{+} = \sum_{K=0,\pm 1} \left(S_{GT,K}^{-} - S_{GT,K}^{+} \right) = 6 \left[\sum_{n_i} v_{n_i}^2 - \sum_{p_i} v_{p_i}^2 \right] = 3 \left(N - Z \right) \tag{59}$$

where we have used Eqs. (56) and (57) to write

$$S_{GT,K}^{-} - S_{GT,K}^{+} = \sum_{\omega_K} \left[\left| \left\langle \omega_K \left| \beta^{-} \right| \phi_0 \right\rangle \right|^2 - \left| \left\langle \omega_K \left| \beta^{+} \right| \phi_0 \right\rangle \right|^2 \right] = \sum_{\gamma_K} \left[\left(b_{\gamma_K} \right)^2 - \left(a_{\gamma_K} \right)^2 \right] \tag{60}$$

It is a simple matter to show that the Ikeda sum rules hold also in the TDA and RPA approximations provided all the ω eigenvalues contained in the basis space are included in the sum so that the orthonormalization conditions

$$\sum_{\omega_K} \left(X_{\gamma_K}^{\omega_K} X_{\gamma'_K}^{\omega_K} - Y_{\gamma_K}^{\omega_K} Y_{\gamma'_K}^{\omega_K} \right) = \delta_{\gamma_K, \gamma'_K} \quad \text{in RPA}$$

and

$$\sum_{\omega_K^{TDA}} X_{\gamma_K}^{\omega_K^{TDA}} X_{\gamma'_K}^{\omega_K^{TDA}} = \delta_{\gamma_K, \gamma'_K} \quad \text{in TDA}$$

are satisfied. In practice, the strength functions are calculated up to an energy cut $\omega \leq E_{cut}$ such that Ikeda's sum rule is satisfied up to a few per thousand.

III. GROUND STATE PROPERTIES AND LOW LYING STRUCTURE

The constrained HF method allows one to get a solution for each value of the mass quadrupole Q_2. In Fig. 1 we show the HF energy as a function of deformation for the two interactions SG2 and Sk3. The best HF solution at each Q_2 value can be obtained by varying the size and deformation parameters [18] of the deformed harmonic oscillator basis. As it is seen in Fig. 1 there are two minima, one in the prolate sector and one in the oblate sector, rather close in energy. Thus, as expected, shape coexistence takes place in this nucleus. This is an interesting feature that was previously predicted (see for instance Refs. [1,24] and references therein) and that has been experimentally confirmed [9]. It should be noted that Fig. 1 leaves open the question as to whether the ground state is oblate or prolate. Sk3 favors prolate shape for the ground state and oblate for the shape isomer, in accordance with previous predictions [25]. On the contrary SG2 force favors an oblate ground state and a prolate shape isomer. The results shown in Fig. 1 exemplify the situation met when several other Skyrme type interactions are used. In particular, the deformation energy curve obtained from calculations performed with the Skyrme force Ska [26] shows a profile quite similar to that of SG2 with an absolute oblate minimum with $Q_2 \simeq -200$ fm^2 and a prolate isomer at about 1 MeV with $Q_2 \simeq 700$ fm ^{2}. On the other hand the force Skm* [27] produces a profile analogous to Sk3 with an absolute prolate minimum at about $Q_2 = 700$ fm^2, and an oblate isomer at $Q_2 = -400$ fm^2. It is also worth mentioning that while the prolate solution appears at about the same deformation for all the forces considered, the oblate solution is spread from $Q_2 = -200$ fm^2 up to $Q_2 = -500$ fm^2 depending on the force considered.

The experimental β_2 values deduced from measured $B(E2)$ strengths [28] vary from $|\beta_2| = 0.24$ to $|\beta_2| = 0.40$. These are consistent with our calculated parameter of deformation β_2 defined in terms of the quadrupole moment Q_2 and the r.m.s. radius $\langle r^2 \rangle$

$$\beta_2 = \sqrt{\frac{\pi}{5}} \frac{Q_2}{A \langle r^2 \rangle} \; . \tag{61}$$

The results obtained for β_2 can be seen in Table 2.

Although in Ref. [9] it is stated that the ground state is prolate and the shape isomer is oblate, we find that experimental evidence supporting this assertion is still lacking and that it would be interesting to investigate further this question from the experimental side. Of course experimental evidence would require measuring the static quadrupole moments of the lowest 2^+ states, however some hints can be obtained from a comparison to data of the theoretical low lying spectra and $E0$ transition rates obtained with these two interactions. To this end we compare in Fig. 2 the experimental [29] level spectra up to 12 MeV with those obtained for each interaction neglecting shape mixing. Clearly, the low lying spectra in Fig. 2 obtained with the prolate Sk3 compare better with experiment. We follow the approximate variation after projection method [30] and write the energy of the $K = 0^+$ band head ($J = 0^+$) as

$$E_{0^+} = E_{HF} - \frac{\langle J^2 \rangle}{2\mathcal{I}_{cr}} \tag{62}$$

and the energies of the higher ($I = 2, 4, 6, ...$) members of the rotational band are calculated to lowest order in angular momentum [31]

$$E_{I^+} - E_{0^+} = \frac{I(I+1)}{2\mathcal{I}_{cr}} \tag{63}$$

where $\mathcal{I}_{cr}$ is the cranking moment of inertia,

$$\mathcal{I}_{cr} = 2 \sum_{\alpha_t \beta_t} \frac{|\langle \alpha_t \beta_t | J_x | \phi_0 \rangle|^2}{E_\alpha + E_\beta} \tag{64}$$

with $|\alpha_t \beta_t\rangle = \alpha^+_{\alpha_t} \alpha^+_{\beta_t} |\phi_0\rangle$ any two quasiparticle state (t = proton, neutron). The values of the Fermi level for neutrons λ_n and protons λ_p, the values of $\langle J^2 \rangle$ and $\mathcal{I}_{cr}$, together with the charge r.m.s. radii, deformations β_2, HF energies (E_{HF}), and 0^+ energies (E_{0^+}, from Eq. (62)), are given in Table 2 for the two minima, oblate and prolate, with the two interactions, SG2 and Sk3. The experimental [20] binding energy is $B(^{74}\text{Kr}) = -631.281$ MeV.

In any case, since the two minima are close in energy and the barriers are not so high, substantial shape mixing can be expected. In Fig. 3 we show that the experimental [9] $\rho^2(E0)$ value is compatible with both Sk3 and SG2 predictions when shape mixing is taken into account. Considering that the $\left|0^+_1\right\rangle$ and $\left|0^+_2\right\rangle$ states mix to form the observed 0^+ ground state $\left|0^+_I\right\rangle$ and first 0^+ excited state $\left|0^+_{II}\right\rangle$

$$|0_I^+\rangle = \sqrt{\lambda}\,|0_1^+\rangle + \sqrt{1-\lambda}\,|0_2^+\rangle \tag{65}$$

$$|0_{II}^+\rangle = \sqrt{1-\lambda}\,|0_1^+\rangle - \sqrt{\lambda}\,|0_2^+\rangle \tag{66}$$

The amplitude of the monopole transition is given by [31]

$$
\begin{aligned}
\rho\left(E0, 0_{II}^+ \to 0_I^+\right) &= \frac{1}{eR^2}\left\langle 0_{II}^+ \left|\hat{E}0\right| 0_I^+\right\rangle \\
&= \frac{1}{eR^2}\left[\sqrt{\lambda(1-\lambda)}\left(\left\langle 0_1^+\left|\hat{E}0\right|0_1^+\right\rangle - \left\langle 0_2^+\left|\hat{E}0\right|0_2^+\right\rangle\right)\right. \\
&\quad \left. - (2\lambda-1)\left\langle 0_2^+\left|\hat{E}0\right|0_1^+\right\rangle\right]
\end{aligned}
\tag{67}
$$

with $\hat{E}0 = \sum_i e_i r_i^2$. The crossed term in Eq. (67) is strictly zero with either of the forces Sk3 or SG2 in the HF limit $\left(v_i^2 = 0, 1\right)$ and can be neglected because of the large β–spacing in between the 0_1^+ and 0_2^+ minima in Fig. 1. One therefore gets for the $\rho^2(E0)$ strength

$$\rho^2(E0) = \lambda(1-\lambda)\left[\frac{Z}{R^2}\left(r_1^2 - r_2^2\right)\right]^2 \tag{68}$$

with r_1 and r_2 the charge r.m.s. radii of the 0_1^+ and 0_2^+ states. In principle, R should be the r.m.s. radius of the ground state but since to our knowledge this has not been measured, we use $R = (r_1 + r_2)/2$ in Eq. (68). An approximate expression for $\rho^2(E0)$ in terms of $\beta_{2(1)}$ and $\beta_{2(2)}$ has been used in Ref. [9,32].

$$\rho^2(E0) = \lambda(1-\lambda)\left[\left(\beta_{2(1)}^2 - \beta_{2(2)}^2\right) + \frac{5\sqrt{5}}{21\sqrt{3}}\left(\beta_{2(1)}^3 - \beta_{2(2)}^3\right)\right]^2 \tag{69}$$

In Fig. 3 we plot $\rho^2(E0)$ as a function of λ, obtained from both Eq. (68) and Eq. (69), using the r.m.s. and deformation values obtained with either Sk3 or SG2 forces (see Table 2). One sees that using Eq. (68) for Sk3 the experimental $E0$ value calls for a large mixing $(\lambda(\lambda-1) \simeq 0.25)$ while for SG2 the mixing must be small $(\lambda(\lambda-1) \simeq 0.05)$. This is consistent with the fact that with SG2 the two minima differ more in energy and the barrier is higher than with Sk3 (see Fig. 1), hence a much smaller mixing probability is to be expected.

We would like to point out that also the $M1$ strength between the lowest 2^+ states can be obtained in an analogous way from the gyromagnetic ratios g_R of the 0_1^+ and 0_2^+ bands given in Table 2. The $B(M1)$ strength between the lowest 2^+ states would be [31]

$$B\left(M1, 2_{II}^+ \to 2_I^+\right) = \frac{648}{4\pi}\lambda(1-\lambda)\left[g_{R,1} - g_{R,2}\right]^2 \mu_N^2 \tag{70}$$

For Sk3 with $\lambda = 0.5$ and $g_{R,1} = 0.487$, $g_{R,2} = 0.459$ (see Table 2), we get $B\left(M1, 2_{II}^+ \to 2_I^+\right) = 0.010~\mu_N^2$ and for SG2 with $\lambda = 0.05$ and $g_{R,1} = 0.498$, $g_{R,2} = 0.463$ (see Table 2), we get $B\left(M1, 2_{II}^+ \to 2_I^+\right) = 0.003~\mu_N^2$.

IV. GAMOW TELLER STRENGTHS AND SUM RULES

The β^+ strength distributions calculated in the selfconsistent HF+RPA scheme described in Section 2 are shown in Fig. 4 as a function of the excitation energy of the daughter nucleus. We present the results for both of the interactions considered, SG2 and Sk3, corresponding to β^+ decay from each of the two minima, oblate and prolate. For comparison we also present the results that would correspond to β^+ decay from the spherical shape. A strong shape dependence is observed in Fig. 4 that is roughly the same independent on whether the SG2 or the Sk3 interaction is used. The Q_{EC} values are indicated by vertical lines. The Q_β value for β^+ decay has been calculated for each of the prolate, oblate, and spherical HF+BCS solutions as

$$Q_{\beta^+} = M_p - M_n - m_e + \lambda_p - \lambda_n - E_p - E_n \tag{71}$$

with λ_p (λ_n) the proton (neutron) Fermi energy given in Table 2 and E_p (E_n) the lowest proton (neutron) quasiparticle energy. The associated Q–value for electron capture is

$$Q_{EC} = Q_{\beta^+} + 2m_e \tag{72}$$

The values of Q_{EC} obtained and shown in Fig. 4 are the following: $Q_{EC} = 4.848$ MeV for the SG2 oblate solution, $Q_{EC} = 4.628$ MeV for the SG2 prolate solution, $Q_{EC} = 4.892$ MeV for the Sk3 oblate solution, and $Q_{EC} = 4.308$ MeV for the Sk3 prolate solution. It is important to stress that the strength summed up to energies below the Q_{EC} window is also strongly dependent on deformation (see also Fig. 8) and this should also be experimentally observable.

Fig. 4 summarizes the main results of this section. One should keep in mind that a quenching of the g_A factor ($g_{A,eff} = 0.7g_A$) is expected on the basis of the observed quenching in charge exchange reactions and spin $M1$ transitions in stable nuclei, where $g_{s,eff}$ is known to be approximately $0.7g_{s,free}$.

In what follows we disentangle the meaning of these results in the light of simpler approximations. In particular, we will study the sensitivity of the Gamow Teller strength distribution to RPA correlations and to pairing correlations, as well as to the nuclear shape.

Since the distinct profiles of the strength distribution in the three nuclear shapes considered are not seriously modified from one interaction to another we restrict the discussion in what follows to results obtained with the SG2 interaction.

In the laboratory frame the transition probability for β^+ decay from the 0^+ to a 1^+_ω state is given by

$$B^+_{GT}\left(0^+ \to 1^+_\omega\right) = \frac{g_A^2}{4\pi}\left|\left\langle 1^+_\omega \left\| \beta^+ \right\| 0^+\right\rangle\right|^2 \tag{73}$$

and the reduced matrix element $\left\langle 1^+_\omega \left\| \beta^+ \right\| 0^+\right\rangle$ in the laboratory frame is related to the intrinsic matrix elements $\left\langle \omega_K \left| \beta^+_K \right| 0\right\rangle$ discussed in the previous section by

$$\left|\left\langle 1^+_\omega \left\| \beta^+ \right\| 0^+\right\rangle\right|^2 = 2\left|\left\langle \omega_1 \left| \beta^+_1 \right| 0\right\rangle\right|^2, \; \omega = \omega_1$$
$$= \left|\left\langle \omega_0 \left| \beta^+_0 \right| 0\right\rangle\right|^2, \; \omega = \omega_0 \tag{74}$$

where we have used the Bohr-Mottelson factorization approximation [31] and have neglected possible higher order corrections due to angular momentum projection [30]. Therefore, the calculated strength functions can be written as

$$B^+_{GT}(\omega) = \frac{g_A^2}{4\pi}\left\{\sum_{\omega_0}\left|\left\langle \omega_0 \left| \beta^+_0 \right| 0\right\rangle\right|^2 \delta\left(\omega - \omega_0\right) + 2\sum_{\omega_1}\left|\left\langle \omega_1 \left| \beta^+_1 \right| 0\right\rangle\right|^2 \delta\left(\omega - \omega_1\right)\right\} \tag{75}$$

consisting of spikes at the various ω_K ($K = 0, 1$) solutions.

To facilitate the comparison among the various approximations discussed here we have folded the calculated GT strengths with $\Gamma = 1$ MeV width Gaussians converting the discrete spectrum into a continuous profile. This can also be interpreted as an approximate way to incorporate the further fragmentation and smoothing effects due to coupling to other excitation modes not taken into account in RPA. The effect of this folding procedure is illustrated in Fig. 5 where we compare the strength in the discrete spectrum to the convoluted strength distribution for the particular case of the prolate solution with the SG2 interaction. In the lower panels the $K = 0$ and $K = 1$ strengths are represented by vertical dotted and solid lines, respectively. The total strength in the upper parts of Fig. 5 is represented by a solid line. The right panels contain the results of the RPA calculations while the left panels contain the results of calculations for β^+ decay to uncorrelated neutron proton two quasiparticle ($2qp$) excitations. The large difference observed between the results of RPA and uncorrelated 2qp neutron-proton excitations in the discrete spectra in Fig. 5 is reliably reflected in the folded distribution. Hence, from here on we restrict our presentation to folded strength distributions.

A. Comparison of RPA, TDA, and bare 2qp strengths

The importance of the role played by the residual interaction is illustrated in Fig. 6 where we compare results for GT strength distributions obtained in RPA (solid line), in TDA (dashed line), and in the uncorrelated two-quasiparticle limit (dotted line). Clearly the repulsive spin-isospin residual interaction (V_{GT}) moves the strength to higher energy. This effect is already present when one goes from the bare 2qp limit to the Tamm Dancoff approximation. The deformation dependence of the strength agrees with that found by Hamamoto and collaborators [12] in the context of TDA. The main effect of the RPA correlations, that can be seen comparing TDA to RPA, is to provide a reduction of the TDA strength maintaining the position of the peaks practically unchanged. A strong dependence of the GT strength on the shape of the β^+ parent is observed in the three approximations, being more dramatic in the

uncorrelated situation. This reflects the fact that the dependence on the different internal shell structure involved in the different shapes is stronger in the absence of a residual interaction. The latter tends to redistribute the strength and tends to produce a smoother strength distribution. This is further illustrated in Fig. 7 where we show the RPA results for oblate, spherical, and prolate shapes corresponding to different coupling strengths of the V_{GT} interaction. The results corresponding to $\chi_{GT} = 0.48$ in Table 1 are compared to those obtained when the strength χ_{GT} of the interaction is reduced or increased by a factor of 2.

The summed strengths up to different excitation energies are represented in Fig. 8 for the two minima (oblate and prolate) corresponding to the SG2 interaction. The results with Sk3 are similar. Also plotted in the figure are the linear energy-weighted sum strengths. Clearly the uncorrelated sum strength

$$\sum_x \left(B_x^{GT\,+}\right)_{uncorr} = \frac{g_A^2}{4\pi} \sum_{\alpha\beta} \left|\langle\alpha\beta\,|\boldsymbol{\sigma}t^+|\,\phi_0\rangle\right|^2 \tag{76}$$

is conserved by the TDA approximation while the RPA approximation conserves the uncorrelated LEWSR

$$\sum_x \left(E_x B_x^{GT\,+}\right)_{uncorr} = \frac{g_A^2}{4\pi} \sum_{\alpha\beta} (E_\alpha + E_\beta) \left|\langle\alpha\beta\,|\boldsymbol{\sigma}t^+|\,\phi_0\rangle\right|^2 \tag{77}$$

where $\alpha\beta$ represents any two quasiparticle neutron-proton excitation. Although this is to be expected from general theory, it is interesting to see that already at low energies ($E^* \lesssim 10$ MeV) in RPA the linear energy-weighted sum tends to converge to the bare 2qp result, while the non energy-weighted sum is clearly lower than the bare 2qp result in order to compensate for the push to higher energies of the strength produced by the V_{GT} interaction. The opposite happens with TDA. The lack of RPA correlations makes the TDA linear energy-weighted sum to start diverging already at 10 MeV. Also shown in Fig. 8 and indicated by arrows, are the calculated Q_{EC} values for the oblate and prolate shapes where one sees that the predicted summed β^+ strength below Q_{EC} is larger in the prolate case. The total linear energy-weighted and non energy-weighted sums up to an excitation energy of 30 MeV are given in Table 3. The results correspond to β^+ decay from the two minima (prolate and oblate) with the SG2 interaction. Also given for comparison are the results for β^+ decay from the prolate minimum with Sk3 interaction. Here one can also see (see Fig. 4) that the results depend more on the shape of the parent nucleus than on the effective interaction. For completeness in Table 3 we give also the β^- summed strength and in the column labelled Ikeda, we give the difference between β^- and β^+ summed strengths to show that the Ikeda sum rule is satisfied in all cases with an accuracy of 0.3%. The summed strengths for $K = 0$ and $K = 1$ are given separately to check that Ikeda's sum rule is satisfied independently for each K value.

B. Analysis of pairing and deformation effects

As already stated our mean field calculations only include neutron-neutron and proton-proton pairing correlations. Taking into account neutron-proton pairing would increase the diffuseness of the Fermi surface (see Goodman [8]), which is governed by the gap parameters. It is therefore interesting to study the sensitivity of the GT strength to this diffuseness. To this end we compare in Fig. 9 the RPA results obtained with the SG2 force for different values of the gap parameters $\Delta_n = \Delta_p = 1.5$ MeV, $\Delta_n = \Delta_p = 1.0$ MeV, and $\Delta_n = \Delta_p = 0.5$ MeV. In the spherical case the role of pairing is different in the low ($E^* \simeq 1$ MeV) and high ($E^* \simeq 6$ MeV) peaks. The first peak decreases with increasing pairing while the second peak increases. In the deformed cases (oblate and prolate) the role of pairing is somewhat less pronounced, but as a general rule, we see that with increasing pairing the strength at high energies (beyond $E^* \simeq 4$ MeV) increases.

To understand better the role of pairing and deformation we show in Fig. 10 the results of bare two quasiparticle excitations with and without pairing for the various shapes. From left to right we can see the effect of going from prolate to spherical and to oblate. From down to upper panels we can see the effect of going from no pairing ($\Delta = 0$) to pairing ($\Delta = 1.5$ MeV). In these figures the solid lines (both in the discrete and continuous distributions) correspond to $K = 1$ excitations and the dotted lines to $K = 0$ excitations. When there is no pairing and no deformation (Fig. 10e) there is only one neutron-proton particle-hole channel open $\left(\pi p_{3/2} \to \nu p_{1/2}\right)$. When we include pairing maintaining the spherical shape (Fig. 10b), the strength of this channel decreases and new channels are open, in particular the $\pi f_{7/2} \to \nu f_{5/2}$ that is forbidden in the case $\Delta = 0, \beta = 0$ (see also Fig. 11). Consequently a two peak structure appears and it remains in the spherical RPA results (see for instance Fig. 4), even though it is modulated by the effect of the residual interaction and RPA correlations. Deformation alone also causes the opening of new channels and the reduction of the $\pi p_{3/2} \to \nu p_{1/2}$ peak (see the transition from panel (e) to panels (d) and (f) in Fig.

10), but in addition it causes fragmentation of the strength between single particle spherical orbitals. In particular, the $K = 0$ and $K = 1$ modes are degenerate in the spherical case independently of whether there is pairing or not. This degeneracy is clearly broken in the deformed case: only $K_\pi^\pm \to K_\nu^\pm$ transitions with $K_\pi = K_\nu = 1/2$ contribute to both $K = 0$ and $K = 1$ modes, while $K_\pi^\pm \to K_\nu^\pm$ transitions with $K_\nu = K_\pi \pm 1$ $(K_\nu, K_\pi > 1/2)$ contribute only to $K = 1$ modes, and those with $K_\nu = K_\pi > 1/2$ contribute only to $K = 0$ modes. Clearly which are the new channels that are open with deformation, and what are their strengths, depends strongly on whether the nucleus is oblate or prolate and on the magnitude of the deformation. As seen in Fig. 10d (see also Fig. 11) in the prolate $\Delta = 0$ case the strongest new channel is the $K = 1$ $(7/2_\pi^- \to 5/2_\nu^-)$, that has a strong spherical component $\pi f_{7/2} \to \nu f_{5/2}$, while the $\pi p_{3/2} \to \nu p_{1/2}$ transition is negligibly small, and other positive parity states have also more strength. On the contrary in the oblate case the strength is still due to the $\pi p_{3/2} \to \nu p_{1/2}$ transition that appears strongly fragmented into several $3/2_\pi^- \to 1/2_\nu^-$ and $1/2_\pi^- \to 1/2_\nu^-$ transitions. The effect of pairing in the deformed case is similar to that in the spherical case. It opens more channels and reduces the strength of the existing ones (compare panels (a) to (d) and (c) to (f) in Fig. 10), resulting in more complex and smoother strength distributions. It is also interesting to observe that in spite of the redistributions of strength caused by the residual interaction V_{GT} and RPA correlations, the distinctive prolate, spherical, and oblate profiles of the B_{GT} strengths in Fig. 4 are still reminiscent of those in Fig. 10. Hence, the above discussion helps to trace back the origin of the main features observed in Fig. 4.

In order to understand better the mechanisms leading to such patterns in the profiles of the GT strength distribution it is illustrative to compare the HF single particle neutron and proton spectra around the Fermi levels in the spherical and deformed cases. This can be seen in Fig. 11, where the Fermi energies are indicated by the step dashed lines. The labels correspond to the ℓ_j values in the spherical case and to K^π in the deformed cases. In the absence of pairing, the allowed GT$^+$ transitions connect only protons below λ_p with neutrons above λ_n with $\Delta K = 0, \pm 1$ and same parity. The most important GT transitions are indicated by the solid arrows. They are $(\pi p_{3/2} \to \nu p_{1/2})$ in the spherical case, $(\pi 7/2^- \to \nu 5/2^-)$ in the prolate case, and $(\pi 1/2^- \to \nu 1/2^-)$ in the oblate one. The dashed arrow indicates the most important GT transition in the spherical case $(\pi f_{7/2} \to \nu f_{5/2})$, which is allowed once pairing correlations are included. From this figure one can also understand why the GT strength in the oblate case is much smaller than in the prolate one. As seen in Fig. 11 the reason for this behaviour is that in the oblate case the proton levels below the Fermi energy are negative parity states while those above the neutron Fermi level are mostly positive parity states, while in the prolate case one finds both, positive and negative parity states, below λ_p and above λ_n. This figure can also be used to predict behaviours of GT strengths when we add or subtract nucleons to ^{74}Kr. For instance, when we remove two neutrons, the state $\nu 3/2^-$ in the oblate case will appear above the Fermi level and therefore it can be fed by β^+ decay of the negative parity proton states below the proton Fermi level. Thus, one expects an important increase of the GT strength when going from the oblate ^{74}Kr to the oblate ^{72}Kr. On the contrary, based on the same picture, we do not expect much changes in the prolate case.

C. Magnetic properties

Gyromagnetic factors and estimates of the total $M1$ transition strength between the lowest 2^+ states were given at the end of Section 3. In this Section we focus on $0^+ \to 1^+$ spin excitations, which are stronger.

The spin $M1$ transitions are the $\Delta T_z = 0$ isospin counterparts of the $\Delta T_z = \pm 1$ Gamow Teller transitions. The study carried out in [13] for several isotope chains showed that the $M1$ strength depends on deformation, not only for orbital but also for spin excitations. Thus, on general grounds, one may argue that the deformation dependence of the GT strength discussed earlier should be comparable to that of the spin $M1$ strength considered in [13].

To see to what extent this argument holds we show in Fig. 12 the profiles of the spin $M1$ strength distributions for the three shapes oblate, spherical, and prolate in ^{74}Kr. The results correspond to selfconsistent HF+RPA calculations with the SG2 interaction. These calculations are described in [13,14] and closely follow the description of calculations in Section 3 except for the $\Delta T_z = 0$ character of the $M1$ operator that, in particular, demands orthogonality of the RPA states to the spurious rotational state [33]. For completeness we also show the results obtained in the absence of residual spin-spin interaction ($K_S = 0$ in Eq. (18)), i.e., the results for bare two quasiparticle spin $M1$ excitations.

One can see in Fig. 12 that as in the case of GT excitations (compare with Fig. 6) the residual interaction produces a redistribution of the strength pushing it to higher energies. There is a striking similarity between the shape dependence in Fig. 12 and that in Fig. 6. This is so in both 2qp and RPA strength distributions.

Although it is unlikely that $M1$ transitions may be observed in such highly unstable nuclei, the comparison of Figs. 6 and 12 tells us that the main features of GT and spin $M1$ strength distributions are very similar. This suggests that we may learn about properties that are observable in highly unstable nuclei (like β decay) from properties that are observable in stable nuclei (like $M1$'s), and vice versa.

V. TOTAL HALF-LIFE

Another quantity of interest in nucleosynthesis is the β^+/EC half-lives of exotic proton rich nuclei. The total half-life depends on the strength distribution in a different way from both the energy-weighted and non-weighted sums considered above. It is therefore illustrative to see the predictions for the half-life as well as the sensitivity of this quantity to the various effects considered in this work.

The partial half-life $t_{1/2}$ for a given transition connecting the ground state of a nucleus with an excited state in the daughter nucleus is given by

$$f\left(Z,\omega\right)t_{1/2} = \frac{D}{\kappa^2 \left|\left\langle 1_\omega^+ \left\|\beta^+\right\| 0^+\right\rangle\right|^2} \tag{78}$$

where we have neglected the Fermi strength. $f\left(Z,\omega\right)$ is the Fermi integral [3,34]. We use $D = 6200$ s and include effective factors

$$\kappa^2 = \left[\left(g_A/g_V\right)_{eff}\right]^2 = \left[0.7\left(g_A/g_V\right)_{free}\right]^2 = 0.74 \tag{79}$$

The total half-life $T_{1/2}$ for allowed β decay from the ground state of the parent nucleus is given by summing over all the final states involved in the process

$$T_{1/2}^{-1} = \frac{0.74}{6200}\sum_\omega f\left(Z,\omega\right)\left|\left\langle 1_\omega^+ \left\|\beta^+\right\| 0^+\right\rangle\right|^2 \tag{80}$$

The Fermi integrals $f\left(Z,\omega\right)$ are taken from Ref. [34].

One should notice that the sum over final states in Eq. (80) is extended over all the states in the daughter nucleus with excitation energies below the $Q-$value. Another difference with respect to the energy-weighted and non-weighted sums considered earlier is that in the sum of Eq. (80) the weighting factor is larger at low excitation energies, which correspond to large kinetic energies of the emitted positron.

We can see in Table 4 the results obtained from bare $2qp$, TDA, and RPA calculations for the total β^+/EC half-life of ^{74}Kr. Results are shown for the two Skyrme forces Sk3 and SG2, as well as for the two shapes oblate and prolate. The experimental value is $\left(T_{1/2}\right)_{EXP} = 11.5$ min. It is clear that the half-lives calculated with both Sk3 and SG2 increase going from $2qp$ to RPA and are larger in the prolate case. On the other hand, the oblate SG2 produces larger $T_{1/2}$ values than the oblate Sk3, while the prolate SG2 gives smaller $T_{1/2}$ values than the corresponding prolate Sk3.

Although not shown in the table we have also checked the sensitivity of this quantity to other effects considered in this work such as the strength of the residual interaction χ_{GT} and the pairing gaps. For example, from an RPA calculation with the force SG2 we find that changing the strength χ_{GT} in Table 1 to half or twice its value, $T_{1/2}$ changes from 3.5 to 16.7 min in the oblate case and from 6.9 to 23.8 min in the prolate case, i.e., roughly speaking, within RPA, $T_{1/2}$ increases proportionally to χ_{GT}. The results are also sensitive to the pairing gaps. Again, an RPA calculation with SG2 and pairing gaps $\Delta_n = \Delta_p = 1$ MeV, gives $T_{1/2} = 1.8$ min and $T_{1/2} = 4.9$ min in the oblate and prolate cases, respectively, i.e., the half-life decreases with decreasing pairing.

One should also keep in mind that values for the effective g_A factors ranging from 0.7 to 0.8 of the bare values are considered in the literature. Since we have taken here the value 0.7, the half-lives in Table 4 would be reduced a 20% by using the value 0.8.

VI. SUMMARY AND FINAL REMARKS

We have investigated shape isomerism and β decay in ^{74}Kr on the basis of the selfconsistent HF+RPA framework with Skyrme forces. This is a well founded method that has been successfully used to describe quite diverse properties of stable spherical and deformed nuclei over the periodic table. The merits of this approximation are well known [13]-[22]. Once the effective interaction parameters are determined by fits to global properties in spherical nuclei over the periodic table and the gap parameters of the pairing force are specified, there are no free parameters left. Both the residual interaction and the mean field are consistently obtained from the same two-body interaction. This is particularly desirable for nuclei far from stability, where extrapolation of methods based on phenomenological mean fields are more doubtful. With the method used here one predicts ground state and low-lying properties, as well as β decay and spin $M1$ excitations without introducing any free parameter. We obtain a fair description of the experimentally known properties of ^{74}Kr and, in particular, of the interesting shape isomerism. In this regard one

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question is open as to whether the ground state is prolate and the first excited 0^+ state is oblate or vice versa, since different Skyrme interactions give different answers. All the Skyrme interactions that we tried (Sk3, SG2, Ska, and Skm*) give a similar stable prolate shape ($\beta_2 \simeq 0.4$), but their predictions on the stable oblate shape are somewhat different ranging from $\beta \simeq -0.15$ up to $\beta \simeq -0.30$. These deformations are compatible with the experimental $B(E2)$ [28] and $\rho^2 (E0)$ [9] values. Theoretical predictions for gyromagnetic ratios and $M1$ transitions are also given.

The RPA Gamow Teller β^+ strength distributions depends strongly on the shape (prolate, spherical or oblate) of the β^+ parent ^{74}Kr, and remarkably, this dependence is quite noticeable in the experimentally accessible energy window. A close similarity is found between this deformation dependence and the deformation dependence of the spin $M1$ strength distribution. It is important to emphasize that these results do not depend much on which effective Skyrme interaction (Sk3 or SG2) is used.

We have also studied the dependence of the results on various aspects of the theory, in particular on the residual interaction, RPA, and pairing correlations. Compared to the uncorrelated two quasiparticle response, RPA shifts the GT strength to higher energies and reduces the total strength. These effects are of course more pronounced if one increases the coupling strength of the residual interaction. While the shifting effect is already contained in the TDA description, the quenching effect is not. This is consistent with the fact that the non energy-weighted summed strength is conserved in TDA, while the linear energy weighted summed strength is conserved in RPA. BCS correlations reduce the strength of allowed particle-hole excitations and create new ones; the main effect of increasing the Fermi diffuseness is to smooth out the profile of the GT β^+ strength distribution increasing the strength at high energies.

Predictions for the half-life are given together with a discussion of the sensitivity of this quantity to the various effects considered in this work. The experimental value, $\left(T_{1/2}\right)_{EXP} = 11.5$ min, is in between the two RPA results for the prolate solutions of the SG2 and Sk3 forces.

It will be interesting to test our results with experimentally accessible information on GT strengths. Since results on the proton rich odd-A Kr isotopes will be soon available [6], it will also be interesting to extend our calculations to include odd-A nuclei. Comparison to data would be desirable before considering further theoretical refinements. The latter would deal with issues such as extended RPA, or the inclusion of the continuum and neutron-proton pairing at the mean field level. It is however important to keep in mind that in the present method the Ikeda sum rule is conserved, while violations of this sum rule are found in some of these extensions.

ACKNOWLEDGMENTS

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Figure 1. Total energy as a function of the mass quadrupole moment obtained from a constraint HF+BCS calculation with the Skyrme interaction SG2 (solid) and Sk3 (dashed).

Figure 2. Experimental spectrum for ^{74}Kr [29] compared to the theoretical rotational spectra obtained for the oblate and prolate solutions of the SG2 and Sk3 interactions.

Figure 3. $E0$ strength from Eq. (68) for SG2 (solid line) and for Sk3 (dotted line). Also shown is the $E0$ strength from Eq. (69) for SG2 (dashed line) and for Sk3 (dash-dotted line). The strength is plotted as a function of the mixing parameter λ (see text). The experimental value from [9] is also indicated.

Figure 4. Gamow Teller strength distribution in ^{74}Kr as a function of the excitation energy of the daughter nucleus. The results correspond to the forces SG2 (solid lines) and Sk3 (dashed lines) in RPA with the coupling strengths given in Table 1. The results are for the prolate (upper part), spherical (middle part), and oblate (lower part) shapes. The vertical solid (dashed) lines indicate the SG2 (Sk3) Q_{EC} values.

Figure 5. Energy distribution of the Gamow Teller strength for the prolate solution of the force SG2 in ^{74}Kr. The lower panels show the spectrum of $K = 0$ (dashed vertical lines) and $K = 1$ (solid vertical lines) excitations. The upper panels show the corresponding folded strength distributions (see text) for $K = 0$ (short dashed line), $K = 1$ (long dashed line), and total (solid line) GT strength. The left (right) panels are from bare two quasiparticle (RPA) calculations.

Figure 6. Comparison of RPA (solid line), TDA (dashed line), and bare two quasiparticle (dotted line) GT strength distributions in ^{74}Kr. The results correspond to the force SG2 for the prolate, spherical, and oblate solutions. Also shown with vertical lines are the Q_{EC} values.

Figure 7. RPA Gamow Teller strength distributions in ^{74}Kr for different coupling strengths of the residual interaction χ_{GT}. The solid lines correspond to the value in Table 1. The short (long) dashed lines correspond to a half (double) value of the coupling strength.

Figure 8. Accumulated GT strength $\left(\Sigma B \left[g_A^2/4\pi \right] \right)$ and energy-weighted strength $\left(\Sigma EB \left[MeV \, g_A^2/4\pi \right] \right)$ in ^{74}Kr as a function of the excitation energy of the daughter nucleus. The results correspond to the prolate (upper part) and oblate (lower part) solutions of SG2. Dotted lines are for bare 2qp calculations, dashed lines for TDA calculations, and solid lines for RPA calculations. The vertical lines indicate the Q_{EC} values.

Figure 9. Pairing effect in the RPA Gamow Teller strength distribution in ^{74}Kr. The solid lines correspond to calculations with pairing gaps $\Delta_n = \Delta_p = 1.5$ MeV, long (short) dashed lines are for $\Delta_n = \Delta_p = 1.0$ MeV ($\Delta_n = \Delta_p = 0.5$ MeV). As in previous figures we can see the results for the prolate (upper part), spherical (middle part), and oblate (lower part) solutions.

Figure 10. Comparison of pairing and deformation effects in the Gamow Teller strength distributions of ^{74}Kr. The results correspond to bare two quasiparticle calculations with the force SG2. From left to right we can see the prolate, spherical, and oblate cases. The upper panels correspond to results with pairing ($\Delta_n = \Delta_p = 1.5 MeV$) and the lower ones are without pairing ($\Delta = 0$). Solid lines (spectra and folded distributions) correspond to $K = 1$ excitations, while dashed lines (spectra and folded distributions) correspond to $K = 0$ excitations. The main configurations leading to the GT excitations are also shown.

Figure 11. Spherical, prolate, and oblate single particle energies for protons and neutrons obtained from HF calculations with the force SG2. The Fermi energies are indicated by the step dashed lines. The most important GT transitions in the spherical and deformed cases are indicated by the solid arrows while the dashed arrow indicates the most important GT transition in the spherical case allowed once pairing correlations are included.

Figure 12. Energy distribution of the spin $M1$ strength with $K^\pi = 1^+$ in ^{74}Kr. The results correspond to the force SG2 and pairing gap parameters $\Delta_n = \Delta_p = 1.5$ MeV. Dashed lines are for 2qp calculations and solid lines for RPA calculations. The three panels are for the prolate (upper part), spherical (middle part), and oblate (lower part) solutions.

Table 1. Parameters of the Skyrme forces SG2 and Sk3: t_0 $\left[\text{MeV fm}^3\right]$, t_1 $\left[\text{MeV fm}^5\right]$, t_2 $\left[\text{MeV fm}^5\right]$, t_3 $\left[\text{MeV fm}^6\right]$, W $\left[\text{MeV fm}^5\right]$, x_0, x_1, x_2, x_3, and α. Also shown are the strengths of the separable isospin χ_F [MeV] and spin-isospin χ_{GT} [MeV] residual interactions obtained from Eqs.(19) and (20) , respectively.

	t_0	t_1	t_2	t_3	W	x_0	x_1	x_2	x_3	$1/\alpha$	χ_F	χ_{GT}
SG2	-2645.0	340.0	-41.9	15595.0	105.0	0.09	-0.0588	1.425	0.06044	6.0	0.69	0.48
Sk3	-1128.75	395.0	-95.0	14000.0	120.0	0.45	0.0	0.0	1.0	1.0	0.88	0.46

Table 2. Fermi energies [MeV] for neutrons (λ_n) and protons (λ_p),values of $\left\langle J^2 \right\rangle$, cranking moments of inertia $\mathcal{I}_{cr}$, gyromagnetic ratios g_R, charge radii r_c, deformation parameters β_2, total energies E_{HF}, and 0^+ energies from Eq. (62) for the two minima (oblate and prolate) of the two interactions, Sk3 and SG2.

		λ_n	λ_p	$\left\langle J^2 \right\rangle$	$\mathcal{I}_{cr}$	g_R	r_c	β_2	E_{HF}	E_{0+}
Sk3	oblate	-12.300	-3.797	21.5	3.6	0.487	4.229	-0.259	-625.40	-628.4
	prolate	-12.589	-4.290	60.3	8.9	0.459	4.269	0.386	-625.89	-629.3
SG2	oblate	-12.739	-3.901	13.4	2.4	0.498	4.147	-0.147	-643.92	-646.7
	prolate	-12.752	-4.139	60.2	9.1	0.463	4.230	0.389	-642.87	-646.2

Table 3. Results obtained from bare 2qp, TDA, and RPA calculations corresponding to the oblate and prolate solutions of SG2 and to the prolate solution of Sk3 in ^{74}Kr. The third (fourth) column contains the GT strength in units of $[g_A^2/4\pi]$ of the β^- (β^+) decay summed up to $E_{cut} = 30$ MeV. The fifth column contains the Ikeda sum rule obtained for this energy cut and the last column contains the energy-weighted summed strength in units of $[MeV\ g_A^2/4\pi]$ corresponding to the β^+ decay. In all cases the results are for $K = 0$ and $K = 1$ (within parentheses).

		$\sum B_{GT}^-$	$\sum B_{GT}^+$	Ikeda	$\sum EB_{GT}^+$
SG2 oblate					
	2qp	4.517 (8.305)	2.523 (4.315)	5.984	12.56 (23.92)
	TDA	4.517 (8.305)	2.523 (4.315)	5.984	16.95 (31.22)
	RPA	3.880 (7.228)	1.886 (3.239)	5.984	12.65 (23.76)
SG2 prolate					
	2qp	5.294 (11.952)	3.300 (7.964)	5.983	17.32 (35.59)
	TDA	5.294 (11.952)	3.300 (7.964)	5.983	23.56 (52.32)
	RPA	4.431 (9.758)	2.437 (5.769)	5.983	17.31 (36.87)
Sk3 prolate					
	2qp	5.389 (12.079)	3.395 (8.088)	5.986	18.86 (37.84)
	TDA	5.389 (12.079)	3.395 (8.088)	5.986	25.28 (54.93)
	RPA	4.533 (9.888)	2.538 (5.898)	5.986	18.79 (39.03)

Table 4. Total half-lives $T_{1/2}$ [minutes] of ^{74}Kr.

	oblate Sk3	oblate SG2	prolate SG2	prolate Sk3
2qp	0.7	1.3	2.8	3.7
TDA	3.3	5.2	7.6	13.9
RPA	4.3	6.9	9.8	16.7

prolate

^{74}Kr

RPA

spherical

oblate

SG2

Sk3

$\mathbf{B(GT)}$ $\mathbf{[g_A{}^2/4\,\pi]}$

Excitation energy [MeV]

prolate

ΣEB

ΣB

oblate

2qp

TDA

RPA

ΣEB

ΣB

Excitation Energy [MeV]

prolate

^{74}Kr

SG 2

RPA

B(GT) [$g_A^2/4\pi$]

spherical

oblate

$\Delta = 1.5$ MeV
$\Delta = 1.0$ MeV
$\Delta = 0.5$ MeV

Excitation energy [MeV]

(a) prolate
ν5/2⁻π7/2⁻ (νf_{5/2}πf_{7/2})
ν1/2⁺π1/2⁺ (νd_{5/2}πg_{9/2})
ν1/2⁺π1/2⁺ (νs_{1/2}πg_{9/2})

(b) spherical
νp_{1/2}πp_{3/2}
νf_{5/2}πf_{7/2}

(c) oblate
K = 0
K = 1

(d)
ν5/2⁻π7/2⁻ (νf_{5/2}πf_{7/2})
ν1/2⁺π1/2⁺ (νd_{5/2}πg_{9/2})
ν1/2⁻π1/2⁻ (νp_{1/2}πp_{3/2})
ν1/2⁺π1/2⁺ (νs_{1/2}πg_{9/2})

(e)
νp_{1/2}πp_{3/2}

(f)
ν1/2⁻π1/2⁻ (νp_{1/2}πf_{5/2})
ν1/2⁻π1/2⁻ (νp_{1/2}πp_{3/2})
ν1/2⁻π1/2⁻ (νp_{1/2}πf_{7/2})

$B(GT)$ [$g_A^2/4\pi$]

Excitation Energy [MeV]

Energy [MeV]

π_{sph} ν_{sph} π_{prol} ν_{prol} π_{obl} ν_{obl}

0

-4

-8

-12

-16

$p_{1/2}$
$f_{5/2}$
$p_{3/2}$

$f_{7/2}$

$g_{9/2}$

$p_{1/2}$
$f_{5/2}$

$p_{3/2}$

$3/2^+$
$3/2^-$
$1/2^-$
$1/2^+$
$7/2^-$

$5/2^-$
$1/2^+$
$1/2^-$
$3/2^-$
$5/2^+$

$3/2^+$

$3/2^-$
$1/2^-$
$1/2^-$
$5/2^-$
$3/2^-$

$1/2^+$
$3/2^+$
$5/2^+$
$7/2^+$
$1/2^-$
$9/2^+$

$3/2^-$